

Information transfer during the universal gravitational decoherence

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Recently Pikovski et al. have proposed in [I. Pikovski, et al. *Nature Phys.* **11**, 668 (2015)] an intriguing universal decoherence mechanism, suggesting that gravitation may play an important role in the quantum-to-classical transition. Here we analyze information transfer induced by this mechanism. We show that generically on the short time-scales, gravitational decoherence leads to a redundant information encoding, governed by energy dispersion and Fisher information. This leads to a objectivization of the center-of-mass position in the gravitational field. As an example we study thermal coherent states and show certain robustness of the effect with the temperature. Finally, we draw an analogy between our objectivization mechanism and the fundamental problem of point individuation in General Relativity as emphasized by the Einstein's Hole argument.

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Emergence of the classical world from quantum has been a long standing problem. Theory of decoherence [1] is one of the attempts to resolve it. In the recent paper [2], Pikovski et al. proposed a universal decoherence mechanism due to the gravitational time dilation [3]. Unlike some of the earlier proposals, e.g. due to Diósi [4] or Penrose [5], the mechanism of Pikovski et al. does not require an immediate departure from unitarity. It has generated a lively debate as to whether the effect is real [6–9] or is it really gravity-related [8, 10]. In this work we will not enter this discussion but rather assume the mechanism of [2] to be true and examine the gravity-mediated information transfer it entails.

Let us recall the idea behind the mechanism of Pikovski et al. [2]. In gravitational field time flows differently depending on the position in the field, which leads to a position-dependent gravitational redshift of frequencies. This effect has been one of the classical tests of general relativity, starting with the experiment of Pound and Rebka [11], and has been recently confirmed over an astonishingly small height difference of 33cm in the Earth gravitational field [12]. Thus, systems which have some natural frequencies associated with their dynamics, like harmonic oscillators, will effectively couple to the position in the gravitational field. If we now promote this reasoning to the quantum domain, then the position-frequency correlation will in general lead to a position decoherence if (some of) the oscillatory degrees of freedom are left unobserved. This, in turn, will lead to a loss of visibility, if one performs an interferometric experiment, and this loss can be directly related to the proper time difference at different heights. Since the gravitational coupling is weak, the effect requires a macroscopic amount of oscillators $N \sim 10^{23}$ to obtain reasonable decoherence rates (milliseconds for micrometric height differences) [2]. A surprising aspect of the mechanism of Pikovski et al. is that there is no position information storage in the oscillator degrees of freedom (which we will simply call the

environment) if the latter are in a thermal equilibrium. Thus, decoherence happens without a localized which-path information and for short times is driven by the internal energy dispersion [2, 13].

Here, to the contrary, we study how the gravitational coupling can deposit position information in the environment if the latter is in a non-stationary state. To study information transfer, we use the approach of the so called Spectrum Broadcast Structures (SBS) [15, 16]—an approach evolved from the quantum Darwinism idea [18, 19] and based on direct studies of an extended quantum state of the system and a part of its environment. Using the simplified model of [2], we show that on a short-timescale gravitational decoherence leads in general to a redundant position information storage in the environment. The efficiency of this storage is governed by the Quantum Fisher Information (QFI) (see e.g. [14]), analogously to decoherence being dependent on the energy dispersion. We study the structure of the extended state and show that it approaches the so called SBS form. As an example, we study in more detail thermal coherent states and the information transfer efficiency as a function of the temperature. Finally, we present an intriguing analogy between the gravity-induced information transfer and the, so called, point individuation in General Relativity (see e.g. [17] for a modern review).

Following [2], we consider a compound system with a large number of internal degrees of freedom, effectively described by uncoupled harmonic oscillators (e.g. a large molecule), placed in a gravitational field. We separate the motion into the center-of-mass (CM) motion and the internal oscillations. With such a division, the CM is treated as a (virtual) central system and the internal oscillators as its environment. In this work, contrary to the previous studies [2, 8, 13], we will assume that only a part of the internal oscillators can be traced out, while the other is left for an observation. This in fact represents a very common situation of an indirect observation [18]

and forces us to study the extended state of the central system plus the observed part of its environment [15, 16]. Assuming weak gravitational field and nonrelativistic velocities, the corresponding approximate Hamiltonian was derived in [2]:

$$H_{tot} = H_{cm}(X, P) + H_0 \left(1 + \frac{\Phi(X)}{c^2} - \frac{P^2}{2m^2c^2} \right), \quad (1)$$

where H_{cm} is a center-of-mass Hamiltonian depending on the canonical variables X, P , H_0 is the internal (oscillatory) degrees of freedom Hamiltonian, and $\Phi(X)$ is the Newtonian gravitational potential. The crucial step is now an assumption that this is a valid Hamiltonian in the quantum domain too, symbolically $H_{tot} \rightarrow \hat{H}_{tot}$. Some further simplifications can be made [2]: i) neglecting H_{cm} as compared to the rest; ii) neglecting the special-relativistic kinematical term $P^2/2M^2c^2$; iii) assuming a homogeneous gravitational field (e.g. from the Earth) $\Phi(X) = gX$, where g is the gravitational acceleration. This finally leads to the Hamiltonian:

$$\hat{H}_{tot} \approx \hat{H}_0 + \hat{H}_{int} = \hat{H}_0 + \Phi(\hat{X}) \otimes \frac{\hat{H}_0}{c^2}, \quad (2)$$

where $\hat{H}_0 \equiv \sum_{i=1}^N \hat{H}_{0i} = \sum_{i=1}^N \hbar \omega_i \hat{n}_i$, $\hat{n}_i = \hat{a}_i^\dagger \hat{a}_i$. This evolution can be easily solved:

$$\hat{U}(t) = \int dX |X\rangle \langle X| \otimes \sum_i \hbar \left(1 + \frac{gX}{c^2} \right) \omega_i \hat{n}_i, \quad (3)$$

where $(1 + gX/c^2)\omega_i \equiv \omega_i(X)$ are the red-shifted frequencies. We apply it to an initial state, assumed to be $\varrho_{tot}(0) = \varrho_0 \otimes \bigotimes_i \varrho_{0i}$, and divide the oscillators into two fractions: the observed one of a size N_o and the unobserved, containing N_\perp oscillators, $N_o + N_\perp = N$. The exact sizes are not important for our considerations, apart from N_o, N_\perp scaling with the total number N (this defines what we call a macrofraction), as we will be tacitly assuming a thermodynamic-type of a limit $N \rightarrow \infty$ [15]. Tracing the unobserved part, gives the extended state:

$$\varrho_{ext}(t) \equiv \text{Tr}_{uno} \varrho_{tot} = \int dX dX' \varrho_0(X, X') \Gamma_t(\Delta X) |X\rangle \langle X'| \otimes \bigotimes_{i \in N_o} U_t^{(i)}(X) \varrho_{0i} U_t^{(i)}(X')^\dagger, \quad (4)$$

where $\varrho_0(X, X') \equiv \langle X' | \varrho_0 | X \rangle$,

$$U_t^{(i)}(X) \equiv \exp[-i\omega_i(X)t\hat{n}_i], \quad (5)$$

and

$$\Gamma_t(\Delta X) \equiv \prod_{i=1}^{N_\perp} \text{Tr} \left[e^{i \frac{g\Delta X \omega_i t}{c^2} \hat{n}_i} \varrho_{0i} \right] \quad (6)$$

is the decoherence factor [2] due to the unobserved oscillators (their number denoted N_\perp). It depends only on the separation $\Delta X \equiv X' - X$.

Let us first perform a short-time analysis [2, 13] for times such that all the phase differences in (6) are small:

$$\Delta\phi_i \equiv \frac{g|\Delta X|}{c^2} \omega_i t \ll 1 \quad (7)$$

or equivalently for times much smaller than the time-dilation induced change of the oscillation periods. Decoherence factor in this regime has been found in [2, 13] and reads:

$$|\Gamma_t(\Delta X)| \approx \exp \left[- \left(\frac{g\Delta X}{\sqrt{2}\hbar c^2} \right)^2 \left(\sum_{i=1}^{N_\perp} \Delta H_{0i}^2 \right) t^2 \right], \quad (8)$$

where $\Delta H_{0i}^2 \equiv \text{Tr}(\varrho_{0i} \hat{H}_{0i}^2) - [\text{Tr}(\varrho_{0i} \hat{H}_{0i})]^2$ is the variance of the i -th oscillator energy, calculated in the initial state ϱ_{0i} . Let us assume that the amount of the unobserved oscillators is very large so that we may use the Law of Large Numbers to write the sum above in a more compact form. This is somewhat similar e.g. to the usual introduction of a spectral density to describe the environment [1]. Let the unobserved frequencies ω be distributed with some probability $p_\perp(\omega)$ such that the average $\langle \langle \Delta H_0^2 \rangle \rangle \equiv \int d\omega p_\perp(\omega) \Delta H_{0\omega}^2$ exists (note that in general the initial states are also ω -dependent). Then:

$$|\Gamma_t(\Delta X)| \approx \exp \left[- \frac{N_\perp g^2 \Delta X^2 \langle \langle \Delta H_0^2 \rangle \rangle}{2\hbar^2 c^4} t^2 \right], \quad (9)$$

giving the decoherence time in a form resembling the energy-time uncertainty relation:

$$\tau_{dec} \sqrt{N_\perp \langle \langle \Delta H_0^2 \rangle \rangle} \equiv \frac{\sqrt{2}\hbar c^2}{g|\Delta X|}. \quad (10)$$

Condition (7) will be fulfilled if for all the relevant frequencies ω the cumulative energy variance is big enough:

$$N_\perp \langle \langle \Delta H_0^2 \rangle \rangle \gg (2\hbar\omega)^2. \quad (11)$$

Let us now study the post-interaction states of the internal oscillators (we assumed that some of them are left for an observation) $\varrho_t^{(i)}(X) \equiv U_t^{(i)}(X) \varrho_{0i} U_t^{(i)}(X)^\dagger$. We are interested in the information they carry about the position X . First of all, since $U_t^{(i)}(X)$ is generated by the red-shifted local Hamiltonian $\hat{H}_{0i}(X)$, it is obvious that any initial state ϱ_{0i} preserved by $\hat{H}_{0i}(X)$, $[\varrho_{0i}, \hat{H}_{0i}(X)] = 0$, will encode no information at all and this is the situation of [2]. However, here we assume a generic ϱ_{0i} . Define $B_t^i(\Delta X) \equiv B[\varrho_t^{(i)}(X), \varrho_t^{(i)}(X')]$, where B is the fidelity (or generalized overlap) $B(\varrho, \sigma) \equiv \text{Tr} \sqrt{\sqrt{\varrho} \sigma \sqrt{\varrho}}$ and is a measure of a state distinguishability [20]. It depends only on the separation as the states $\varrho_t^{(i)}(X)$ differ for different X only by an unitary rotation depending on $\Delta\phi_i$ (7). For short times (7), we develop B in a series in $\Delta\phi$, using first the definition of the Bures distance d_B [21] $B(\Delta\phi) = 1 - 1/2[d_B(\Delta\phi)]^2$ and then the fact

that an infinitesimal Bures distance is given by 1/4 of the quantum Fischer information (QFI) $F(\varrho_{0i}; \hat{H}_{0i})$ [22]:

$$B_t^i(\Delta X) \approx 1 - \frac{t^2}{8} \left(\frac{g\Delta X}{\hbar c^2} \right)^2 F(\varrho_{0i}; \hat{H}_{0i}). \quad (12)$$

By the quantum Cramer-Rao bound, $F(\varrho_{0i}; \hat{H}_{0i})$ sets the limit on the estimation of $gt\Delta X/(\hbar c^2)$. It can be given in terms of the eigenvalues and eigenvectors of the initial state $\varrho_{0i} = \sum_n \lambda_n |\lambda_n\rangle \langle \lambda_n|$ as [14]:

$$F(\varrho_{0i}; \hat{H}_{0i}) = 2 \sum_{m,n} \frac{(\lambda_n - \lambda_m)^2}{\lambda_n + \lambda_m} |\langle \lambda_n | \hat{H}_{0i} | \lambda_m \rangle|^2. \quad (13)$$

If ϱ_{0i} is diagonal in the basis of \hat{H}_{0i} then obviously $F(\varrho_{0i}; \hat{H}_{0i}) = 0$ and there is no encoding as mentioned earlier. Let us now further divide the observed fraction into a number of smaller macrofractions mac of a size $N_{mac} = \mu N_o$, $0 < \mu < 1$, large enough to use the LLN [23]. They may represent e.g. portions of the environment accessible to multiple observers [15, 16]. Crucially, the fidelity separates w.r.t. the tensor product $B(\bigotimes \varrho_i, \bigotimes \sigma_i) = \prod B(\varrho_i, \sigma_i)$, so we can use the same analysis as before, applied to the macrofraction states $\varrho_t^{mac}(X) \equiv \bigotimes_{i \in mac} \varrho_t^{(i)}(X)$ and their fidelities. Introduce a probability $p_{mac}(\omega)$ such that the average QFI $\langle \langle F_0 \rangle \rangle \equiv \int d\omega p_{mac}(\omega) F[\varrho_0(\omega); \hat{H}_{0\omega}]$ exists. We obtain a formula similar to (9) but with the average variance substituted for the average quantum Fischer information:

$$B_t^{mac}(\Delta X) \approx \exp \left[- \frac{N_{mac} g^2 \Delta X^2 \langle \langle F_0 \rangle \rangle}{8 \hbar^2 c^4} t^2 \right]. \quad (14)$$

This defines the distinguishability time-scale in analogy to the decoherence time-scale (10):

$$\tau_{dst} \sqrt{N_{mac} \langle \langle F_0 \rangle \rangle} \equiv \frac{\sqrt{8} \hbar c^2}{g |\Delta X|}. \quad (15)$$

Condition (7) is fulfilled when the cumulative QFI is large compared to the relevant energies:

$$N_{mac} \langle \langle F_0 \rangle \rangle \gg (8 \hbar \omega)^2. \quad (16)$$

We note that since $F(\varrho_{0i}; \hat{H}_{0i}) \leq 4 \Delta H_{0i}^2$ [14], $\tau_{dst} \geq \tau_{dec}$ provided the observed and the unobserved macrofractions are identical.

The decoherence and distinguishability times strongly depend on the separation $|\Delta X|$. One can view (9, 14) from another perspective: For a fixed time t satisfying (7), there are characteristic decoherence and distinguishability lengths:

$$\Delta X_c^2 \equiv \frac{2 \hbar^2 c^4}{g^2 t^2 N_{\perp} \langle \langle \Delta H_0^2 \rangle \rangle}, \quad (17)$$

$$\Delta X_d^2 \equiv \frac{8 \hbar^2 c^4}{g^2 t^2 N_{mac} \langle \langle F_0 \rangle \rangle} \quad (18)$$

such that for $|\Delta X| > \Delta X_c$, we may assume $\Gamma_t(\Delta X) \approx 0$ and for $|\Delta X| > \Delta X_d$, $B_t(\Delta X) \approx 0$. Let us study the consequences of this for the state (4). We fix t , and change the integration variables to $\bar{X} \equiv 1/2(X + X')$ and ΔX . Then, since $\Gamma_t(\Delta X) \approx 0$ for $|\Delta X| > \Delta X_c$, we have:

$$\begin{aligned} \varrho_{ext}(t) &\approx 2 \int d\bar{X} \int_{-\Delta X_c}^{\Delta X_c} d\Delta X \varrho_0(\bar{X} - \Delta X, \bar{X} + \Delta X) \\ &\Gamma_t(\Delta X) |\bar{X} - \Delta X\rangle \langle \bar{X} + \Delta X| \otimes \\ &\bigotimes_{i \in N_o} U_t^{(i)}(\bar{X} - \Delta X) \varrho_{0i} U_t^{(i)}(\bar{X} + \Delta X)^\dagger. \end{aligned} \quad (19)$$

We now break the first integral into intervals of length ΔX_d , centered at some points \bar{X}_k : $\int d\bar{X} f(\bar{X}) = \sum_k \int_{-\Delta X_d}^{\Delta X_d} d\Delta \bar{X} f(\bar{X}_k + \Delta \bar{X})$. We then expand everything that is smooth in $\Delta X, \Delta \bar{X}$ and keeping the lowest term obtain:

$$\begin{aligned} \varrho_{ext}(t) &= 2 \sum_k p_0(\bar{X}_k) \int_{-\Delta X_c}^{\Delta X_c} d\Delta X \int_{-\Delta X_d}^{\Delta X_d} d\Delta \bar{X} \\ &|\bar{X}_k + \Delta \bar{X} - \Delta X\rangle \langle \bar{X}_k + \Delta \bar{X} + \Delta X| \otimes \bigotimes_{i \in N_o} \varrho_t^{(i)}(\bar{X}_k) \\ &+ O(\Delta X_c^2, \Delta X_d^2), \end{aligned} \quad (20)$$

where $p_0(X) = \varrho_0(X, X) = \langle X | \varrho_0 | X \rangle$. Changing once again the integration variables to sum and difference of $\Delta X, \Delta \bar{X}$ and defining smeared position states $|\bar{X}; \bar{\Delta}\rangle \equiv \int_{-\bar{\Delta}}^{\bar{\Delta}} d\Delta X |X + \Delta X\rangle$, where $\bar{\Delta} \equiv (\Delta X_c + \Delta X_d)/2$, we finally obtain:

$$\begin{aligned} \varrho_{ext}(t) &\approx \\ &\sum_k p_0(\bar{X}_k) |\bar{X}_k; \bar{\Delta}\rangle \langle \bar{X}_k; \bar{\Delta}| \otimes \varrho_t^{mac}(\bar{X}_k) \dots \varrho_t^{mac}(\bar{X}_k). \end{aligned} \quad (21)$$

The macrofraction states $\varrho_t^{mac}(\bar{X}_k)$ are by the above construction almost perfectly distinguishable for different \bar{X}_k , since their separation is $\sim \Delta X_d$. Thus, the CM position is decohered to within $\bar{\Delta}$ and is stored in the internal degrees of freedom in many identical copies. The structure (21) is a version of the Spectrum Broadcast Structure for continuous variables. The question of the approach to SBS was analyzed in [25].

As an illustration of the above general reasoning, we consider the internal degrees of freedom to be prepared in thermal coherent states:

$$\varrho_{0i} = \hat{D}(\alpha) \varrho_{th}^{(i)} \hat{D}(\alpha)^\dagger, \quad (22)$$

where $\hat{D}(\alpha)$ is the displacement operator (assumed here the same for all the oscillators) and $\varrho_{th}^{(i)} \equiv e^{-\beta \hat{H}_{0i}} / Z_i$, $Z_i \equiv \text{Tr}(e^{-\beta \hat{H}_{0i}})$, $\beta \equiv 1/(kT)$. Those states obviously break the symmetry w.r.t. \hat{H}_{0i} and we have:

$$\Delta H_{0i}^2 = (\Delta H_{0i}^2)_{th} + (\hbar \omega_i |\alpha|)^2 \text{cth} \left(\frac{\beta \hbar \omega_i}{2} \right), \quad (23)$$

$$F(\varrho_{0i}; \hat{H}_{0i}) = 4(\hbar \omega_i |\alpha|)^2 \text{th} \left(\frac{\beta \hbar \omega_i}{2} \right), \quad (24)$$

where $\langle \hat{A} \rangle_{th} \equiv \text{Tr}(\varrho_{th} \hat{A})$. We note the reciprocal dependence on the temperature of the non-thermal parts above, which formally leads to a type of an information-vs.-disturbance relation (cf. e.g. [24]):

$$[\Delta H_{0i}^2 - (\Delta H_{0i}^2)_{th}] F(\varrho_{0i}; \hat{H}_{0i}) = 4(\hbar\omega_i |\alpha|)^4. \quad (25)$$

Here, the disturbance to the central system (in a form of a non-thermal contribution to the decoherence) is characterized by $\Delta H_{0i}^2 - (\Delta H_{0i}^2)_{th}$, while the environment information gain by $F(\varrho_{0i}; \hat{H}_{0i})$. The meaning of (25) here is that the hotter the environment is, the stronger it decoheres but the lesser is its information capacity [26].

One can go beyond the above short-time analysis and give the exact $|\Gamma_t(\Delta X)|$ and $B_t(\Delta X)$ for arbitrary times. A direct calculation for a single band ω_i leads to:

$$|\Gamma_t^i(\Delta X)| = |\Gamma_t^{th}(\Delta X)| \times \exp \left[-|\alpha|^2 |\Gamma_t^{th}(\Delta X)|^2 \text{cth} \left(\frac{\beta \hbar \omega_i}{2} \right) (1 - \cos \Delta \phi_i) \right], \quad (26)$$

where $|\Gamma_t^{th}(\Delta X)| \equiv [1 + 2\bar{n}_i(\bar{n}_i + 1)(1 - \cos \Delta \phi_i)]^{-1/2}$, $\bar{n}_i \equiv \langle \hat{n}_i \rangle_{th}$. Calculation of $B_t^i(\Delta X)$ is more involved. Using the methods of [26], we find:

$$B_t^i(\Delta X) = \exp \left[-|\alpha|^2 \text{th} \left(\frac{\beta \hbar \omega_i}{2} \right) (1 - \cos \Delta \phi_i) \right]. \quad (27)$$

The above single-band functions are periodic in time (hidden in $\Delta \phi_i$ (7)) and obviously there is no SBS formation at the level of single environments. However, for macrofractions, described by a collection of randomly distributed ω_i , the corresponding decoherence and fidelity factors become quasi-periodic functions of time and the random phases $\Delta \phi_i$ may lead to their effective damping. This, however, depends on the temperature (apart from the other factors kept fixed) as e.g. for a high temperature $\frac{k_B T}{\hbar \omega_i} \rightarrow \infty$, $|\Gamma_t^i(\Delta X)| = O\left(\frac{\hbar \omega_i}{k_B T}\right)$ while $B_t^i(\Delta X) = 1 - O\left(\frac{\hbar \omega_i}{k_B T}\right)$. For a low temperature in turn, the initial states (22) become pure and $B_t^i(\Delta X) = |\Gamma_t^i(\Delta X)|$, which follows from $B(|\psi\rangle\langle\psi|, |\phi\rangle\langle\phi|) = |\langle\psi|\phi\rangle|$. An example of the intermediate regime $k_B T \approx \hbar \omega_i$ is shown in Fig. 1.

We finish by drawing some analogy between the above information transfer process and the, so called, point individuation in General Relativity. As emphasized by the famous Einstein's Hole argument, active diffeomorphism invariance of General Relativity forbids assigning a physical meaning not only to coordinate charts but to very (mathematical) manifold points as well. This has led to a still active debate over the ontology of space-time and physical objectivity of space-time points [17, 27]. A way out of the situation is to define space-time points using coincidences of matter, e.g. as intersection points of particles' world-lines. This process is called space-time point individuation or objectivization and the used degrees of

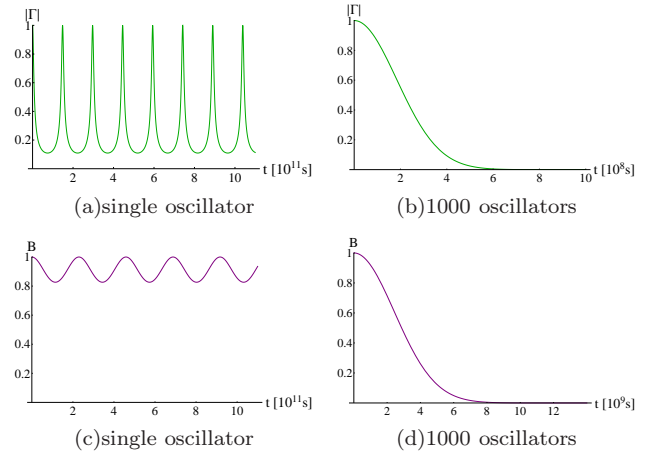


FIG. 1: (Color online) Time dependencies of $|\Gamma_t(\Delta X)|$ (a),(b) and $B_t(\Delta X)$ (c),(d) for intermediate temperatures. Plots (a),(c) show a single-oscillator case with $\omega = 2.5 \times 10^{11} \text{s}^{-1}$; (b),(d) show $N_{\perp} = N_{mac} = 10^3$ random oscillators, uniformly distributed over $[1 \dots 5] \times 10^{11} \text{s}^{-1}$; $\Delta X = 10^{-6} \text{m}$, $|\alpha| = 1$, $T = 10 \text{K}$ ($k_B T \approx 3\hbar\omega_{max}$). Note the extremely long decay times in (b),(d) due to the small macrofraction sizes as compared to the gravitational interaction strength.

freedom-individuating field. The peculiarity of the theory is that individuation can be in principle achieved in a fully dynamical way by the metric field itself, provided it satisfies the Einstein equations [28].

On the other hand, the SBS structure (21) encodes a certain form of point objectivity as well: Indirect observation of the position through any of the environmental macrofractions will always give the same result \bar{X}_k , leaving the CM located around \bar{X}_k and without any disturbance, on average, to the global state. This observer-invariance and non-disturbance (which may be viewed as a type of a time-invariance) can be taken as a basis of an operational definition of objectivity [18], in this case of the (approximate) position. Moreover, it can be shown that under some general assumptions the only state structure compatible with so defined objectivity is precisely the SBS [16] and that it is generic for macroscopic quantum measurements [29]. An intriguing question thus arises if the gravitational decoherence mechanism may also lead, through the SBS, to the point individuation? This would be an example of a quantum individuation.

There are of course some obvious differences. First of all, described here gravity-generated SBS provides objectivization of points of space only, not of space-time events (see however [5]). But we have used a non-relativistic Schrödinger equation rather than e.g. Tomonaga-Schwinger evolution law, which could be a future direction. Second, by the very construction there is a finite precision (17) and resolution (18) with which a point can be localized. Both of these parameters improve with the macrofraction size, which sug-

gests that objective points may be macroscopic phenomena and the objectivity breaks down at some microscopic scales (cf. macro-objectivity idea [30]). Answers to these questions may shed some light on what happens with the space-time at microscopic scales.

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